PUTP-94-10

The Quark Axial Vector Coupling and Heavy Meson Decays

Nan-Guang Chen

Department of Physics, Peking University, Beijing 100871, P.R.China

and

Kuang-Ta Chao

CCAST (World Laboratory), Beijing 100080, P.R.China

and

Department of Physics, Peking University, Beijing 100871, P.R.China

Abstract

Form factors and decay widths for $D^* \to D\gamma$ and $D^* \to D\pi$ decays are estimated in a relativistic constituent quark model. Relativistic corrections due to light quarks are found to be substantial and to suppress the vector and axial vector form factors. The $CLEO$ experimental value of $R^0_\gamma \equiv \Gamma(D^{*0} \to D^0\gamma)/\Gamma(D^{*0} \to D^0\pi^0) = 0.572 \pm 0.057 \pm 0.081$ is used to determine the quark axial vector coupling g_A, which is found to be $0.6-0.8$ for $m_u = (350-200)MeV$ correspondingly, as compared with the chiral model result $g_A = 0.8 - 0.9$. The heavy meson-pion strong coupling g is found to be $0.4-0.6$, much smaller than $g = 1$ which is expected in the large N_C and nonrelativistic limit, but consistent with some heavy hadron chiral theory and QCD sum rule results.

Because the chiral $SU(2)$ symmetry is spontaneously breaking, the axial vector coupling of the constituent quark, g_A, may take any values. It has recently been argued based on a chiral model[1] that to leading order in $1/N_C$ (where N_C is the number of colors in quantum chromodynamics), the constituent quarks just behave like bare Dirac particles, and both u and d quarks have vanishing anomalous magnetic moment and have the axial vector coupling

$$g_A = 1. \tag{1}$$

With order $1/N_C$ corrections, the value of g_A is slightly reduced and is estimated to be[2,3]

$$g_A = 0.8 - 0.9. \tag{2}$$

On the other hand, using the experimental value for the nucleon axial vector coupling $G_A = 1.25$ with the static SU(6) wave function for the nucleon leads to $g_A = 1.25 \times \frac{3}{5} = 0.75$. The discrepancy between this value and (2) might be removed by including the relativistic corrections to the nucleon wave functions[4].

The heavy meson decays, e.g., $D^* \to D$ decays may also be a good testing ground for the value of quark axial vector coupling g_A. Although the total widths of D^{*+} and D^{*0} are still unknown, their decay branching ratios have been measured accurately by the $CLEO$ Collaboration[5] (see also ref.[6]). These heavy meson decays have been discussed in the literature[7−13]. To determine g_A, the most useful data might be $B(D^{*0} \to D^0\gamma)$ and $B(D^{*0} \to D^0\pi^0)$, which are better measured experimentally and are not sensitive to the value of the charm quark mass in theoretical calculations. We suggest using the measured ratio[5]

$$R^0_\gamma \equiv \frac{\Gamma(D^{*0} \to D^0\gamma)}{\Gamma(D^{*0} \to D^0\pi^0)} = 0.572 \pm 0.057 \pm 0.081, \tag{3}$$

together with the calculated $\Gamma(D^{*0} \to D^0\gamma)$ and the form factors in the hadronic decay to determine g_A.

Because the light constituent quarks (u and d) inside the heavy mesons are relativistic, to discuss the g_A issue in the heavy meson decays the relativistic motion of light quarks has to be taken into consideration. In the following we will study the heavy meson decays and the g_A problem in a relativistic constituent quark model, based on the Bethe-Salpeter (BS) formalism.

For a $q_1\bar{q}_2$ bound system with quark momentum q_1, mass m_1; antiquark momentum q_2, mass m_2; relative momentum q; and total momentum P, meson mass M; and

$$p_1 = \eta_1 P + q,\ p_2 = \eta_2 P - q,\ \eta_i = \frac{m_i}{m_1 + m_2}\ (i = 1, 2),$$

the general form of the three dimensional BS wave functions for the 0^- and 1^- mesons

is given by (see e.g. ref.[14])

$$\Phi_P^P(\vec{q}) = \Lambda_+^1(\vec{p_1})\gamma^0(1+\frac{\hat{P}}{M})\gamma_5\gamma^0\Lambda_-^2(\vec{p_2})\phi(\vec{q}) \quad (4)$$

$$\Phi_P^V(\vec{q}) = \Lambda_+^1(\vec{p_1})\gamma^0(1+\frac{\hat{P}}{M})\hat{e}\gamma^0\Lambda_-^2(\vec{p_2})f(\vec{q}) \quad (5)$$

where e_μ is the polarization vector of the 1^- meson, $e_\mu P^\mu = 0; E_i = \sqrt{\vec{p_i}^2 + m_i^2}$, and $\Lambda_\pm^i = \frac{1}{2}(1 \pm \frac{\gamma^0(\vec{\gamma}\cdot\vec{p_i}+m_i)}{E_i})$ are the positive (+) and negative (-) energy projectors, and ϕ, f are scalar wave functions. (4) and (5) respect the space reflection symmetry, and also respect the flavor-spin symmetry in the heavy quark limit. This can be seen by taking $m_1 \to \infty$, then $p_1^\mu \to P^\mu$ and (4) and (5) become

$$\Phi_P^P(\vec{q}) = \frac{1}{v^0}(1+\hat{v})\gamma_5\gamma^0\Lambda_-^2(\vec{p_2})\phi(\vec{q}), \quad (6)$$

$$\Phi_P^V(\vec{q}) = \frac{1}{v^0}(1+\hat{v})\hat{e}\gamma^0\Lambda_-^2(\vec{p_2})f(\vec{q}), \quad (7)$$

where $v^\mu = \frac{P^\mu}{M}$, and $f = \phi$ which is due to the vanishing of the color magnetic force in the heavy quark limit. The normalization condition for (4) and (5) is

$$\frac{1}{2\pi}\int d^3qTr\left[\Phi_P^\dagger(\vec{q})\Phi_P(\vec{q})\right] = 2P^0 = 2\sqrt{\vec{P}^2+M^2}. \quad (8)$$

Suppose a flavor changing quark operator $\bar{q}_f\Gamma q_i$ induces the transition

$$\Phi_P(\vec{q})(p_1, m1; p_2, m_2; P, M) \to \Phi_{P'}(\vec{q'})(p_1', m_1'; p_2, m_2; P', M'), \quad (9)$$

where the antiquark (p_2, m_2) remains a spectator, then the transition matrix element is given by

$$\langle P' \mid \bar{q}_f\Gamma q_i \mid P\rangle = \frac{1}{2\pi}\int d^3p_2Tr\left[\Phi_{P'}^\dagger(\vec{q'})\gamma^0\Gamma\Phi_P(\vec{q})\right]. \quad (10)$$

Likewise, the matrix element induced by the antiquark transition can be written in a similar way.

In the radiative $M1$ decay such as $D^* \to D\gamma$ if we neglect the quark anomolous magnetic moment[1,2] we then only need to calculate the matrix element induced by the vector operator $\bar{q}\gamma_\mu q$ of the quark and antiquark, and according to (10), (4), and (5), for the quark transition we have

$$\begin{aligned} j_\mu &= \langle P' \mid \bar{q_1}\gamma_\mu q_1 \mid P, e\rangle \quad (11)\\ &= \frac{1}{2\pi}\int d^3p_2Tr\left[\frac{\hat{p_2}-m_2}{2E_2}\gamma_5(1+\hat{v'})\frac{\hat{p_1'}+m_1}{2E_1'}\gamma_\mu\frac{\hat{p_1}+m_1}{2E_1}(1+\hat{v})\hat{e}\right]\\ &\quad \times f_i(\vec{p_2}-\frac{m_2}{m_1+m_2}\vec{P})\phi_f(\vec{p_2}-\frac{m_2}{m_1+m_2}\vec{P'}), \end{aligned}$$

where $\hat{v} = \frac{\hat{P}}{M_i}, \hat{v}' = \frac{\hat{P}'}{M_f}$. In general, the vector current matrix element (11) can be expressed as

$$j_\mu = -i\sqrt{M_i M_f}\frac{M_f}{m_1}\epsilon_{\mu\nu\alpha\beta}e^\nu v^\alpha v'^\beta \xi_{V1}, \tag{12}$$

where ξ_{V1} is the vector form factor due to the quark q_1 transition. If the scalar wave functions f_i (for the initial state 1^- meson) and ϕ_f (for the final state 0^- meson) in (11), which are to be determined by the interquark forces, are known, then a direct calculation for (11) will give the form factor ξ_{V1}. We will give this result later on by solving the BS equation with a QCD-motivated interquark potential. Before doing that calculation we may first consider a simplified calculation as follows. Because in the heavy quark limit the color-magnetic force vanishes, the 0^- and 1^- mesons will have the same spatial wave functions, we may assume f_i and ϕ_f to take the same form

$$f_i(\vec{p}) = a_i f(\vec{p}), \qquad \phi_f(\vec{p}) = a_f f(\vec{p}), \tag{13}$$

where a_i and a_f are normalization factors determined by (8) in the initial and final meson frames respectively. Although in (11) the light quark is relativistic, in order to see the relativistic effects more explicitly (but less rigorously) it might be instructive to make a nonrelativistic expansion in terms of the inverse of the quark masses. Then to the first order we find

$$\xi_{V1} = 1 - (\frac{2}{3m_1^2} - \frac{1}{8m_1 m_2})\langle\vec{q}^2\rangle - \frac{1}{6}(\frac{m_2}{m_1+m_2})^2 \mid \vec{K}\mid^2 \langle\vec{r}^2\rangle, \tag{14}$$

where $\vec{K} = \vec{P}' - \vec{P}$ is the recoil momentum, $\langle\vec{q}^2\rangle = \int d^3q \mid f(\vec{q})\mid^2 \vec{q}^2$ is the mean value of the quark momentum squared, and $\langle\vec{r}^2\rangle$ is the mean value of the radius squared of the mesons. From (14) we may find some qualitative feature of the relativistic effects. We see that the M1 transition can be substantially suppressed if the quark is a light quark and becomes relativistic $\frac{\langle\vec{q}^2\rangle}{m_1^2} = O(1)$. We see also that the effect of the spectator antiquark on the quark transition is not strong in general, because the coefficient of the term involving $(m_1 m_2)^{-1}$ is small. The last term in (14) is due to the nonzero recoil momentum, and its contribution is also small because of the smallness of the recoil momentum in the D^* radiative decays. If the quark is a heavy quark then in the heavy quark limit $m_1 \to \infty$ we will have $\xi_{V1} = 1$ in (14) and $\frac{M_f}{m_1} \to 1$ and therefore (12) will return to the well known expression for the vector current in the heavy quark limit.

In the rest frame of the 1^- meson, $\vec{v} = 0, v^0 = 1, \vec{v'} = \frac{-\vec{K}}{M_f}$, then (12) gives the familiar expression for the M1 transition amplitude. With a similar expression for the antiquark transition, we can get the $1^- \to 0^-\gamma$ decay width for a $q\bar{Q}$ system

$$\Gamma_\gamma = \frac{\alpha}{3}(\xi_q \frac{e_q}{m_q} + \xi_Q \frac{e_Q}{m_Q})^2 \mid \vec{K}\mid^3, \tag{15}$$

where ξ_q and ξ_Q are the vector form factors ξ_{V1} for $q = u,\ d$ and $Q = c$ in the $D^* \to D\gamma$ decay. We will calculate these form factors from (11) and will not use their nonrelativistic expansion (14), because the light quark can be highly relativistic.

The hadronic decay $D^* \to D\pi$ can be described by the quark-pion vertex in the chiral quark model[15]

$$\mathcal{L}_I = -\frac{g_A}{2\sqrt{2}F_\pi}\bar{\psi}\gamma^\mu\gamma_5\tau_i\psi\partial_\mu\pi^i, \tag{16}$$

where g_A is the quark axial vector coupling and $F_\pi = 132MeV$ is the pion decay constant. In general, for a $q_1\bar{q}_2$ system the transition induced by the quark axial vector current can be written as

$$\begin{aligned} j_5^\mu &= \langle P' \mid \bar{q}_1\gamma^\mu\gamma_5 q_1 \mid P, e\rangle \\ &= \frac{1}{2\pi}\int d^3p_2 Tr\left[\frac{\hat{p_2}-m_2}{2E_2}\gamma_5(1+\hat{v}')\frac{\hat{p_1'}+m_1}{2E_1'}\gamma^\mu\gamma_5\frac{\hat{p_1}+m_1}{2E_1}(1+\hat{v})\hat{e}\right] \\ &\quad \times f_i(\vec{p_2}-\frac{m_2}{m_1+m_2}\vec{P})\phi_f(\vec{p_2}-\frac{m_2}{m_1+m_2}\vec{P'}) \\ &= \sqrt{M_{D^{*0}}M_{D^0}}\left\{\xi_{A1}(1+v\cdot v')e^\mu - \xi_{A2}(e\cdot v')v^\mu - \xi_{A3}(e\cdot v')v'^\mu\right\}. \end{aligned} \tag{17}$$

As in the case of ξ_{V1}, we can calculate ξ_{A1}, ξ_{A2}, and ξ_{A3} with the scalar wave functions f and ϕ obtained by solving the BS equation, but it is useful to give the nonrelativistic reduction form for, e.g., the ξ_{A1}

$$\xi_{A1} = 1 - \frac{\langle\vec{q}^2\rangle}{3m_1^2} - \frac{1}{6}(\frac{m_2}{m_1+m_2})^2 \mid \vec{K}\mid^2 \langle\vec{r}^2\rangle. \tag{18}$$

Again, ξ_{A1} is suppressed by the relativistic motion of the light quark, but the suppression is less severe than the M1 transition form factor ξ_{V1} in (14). This can be seen by noting that the coefficient of the $\frac{\langle\vec{q}^2\rangle}{m_1^2}$ term, which is the leading term for the suppression, is $-1/3$ in (18) whereas is $-2/3$ in (14).

For e. g. the $D^{*0} \to D^0\pi^0$ decay, $\mid \vec{P_\pi} \mid = 44MeV, \mid \vec{v'} \mid = 0.024, \mid e\cdot v' \mid \ll 1$, therefore in (17) the contribution of ξ_{A2} and ξ_{A3} terms can be neglected, and we then get

$$\Gamma(D^{*0}\to D^0\pi^0) = \frac{g_A^2\xi_{A1}^2}{12\pi F_\pi^2} \mid \vec{P_\pi}\mid^3. \tag{19}$$

For $D^{*+}\to D^+\pi^0$ the width takes the same form as (19), while for $D^{*+}\to D^0\pi^+$ an additional factor of 2 on the right hand side is needed due to the isospin difference.

To calculate the form factors in radiative decay (11) and (15) and in hadronic decay (17) and (19), as a simple and naive choice, we first use the Gaussian wave functions

$$f(\vec{q}) = Ne^{-\vec{q}^2/a^2} \tag{20}$$

for f_i and ϕ_f, where N is the normalization factor, $a^2 = \frac{4}{3}\langle \vec{q}^2 \rangle$. For the D and D^* mesons most estimates give (see, e.g., ref.[16])

$$\langle \vec{q}^2 \rangle = 0.2 - 0.3 \; GeV^2 \tag{21}$$

Here we will simply take $\langle \vec{q}^2 \rangle = 0.21 \; (0.30) \; GeV^2$ for $m_u = m_d = 200 \; (300) \; MeV$, and $m_c = 1500 \; MeV$, and calculate the form factors using their relativistic expressions (11) and (17) (not (14) and (18)). The calculated form factors ξ_u, ξ_c, ξ_{A1}, and the decay widths are shown in Table 1. We see these numerical results are qualitatively consistent with the nonrelativistic reduction expressions (14) and (18) (e.g., ξ_u is more suppressed than ξ_{A1}). Here g_A is determined by using the calculated width for $D^{*0} \to D^0\gamma$ and the experimental value for the ratio $R_\gamma^0 \equiv \frac{\Gamma(D^{*0} \to D^0\gamma)}{\Gamma(D^{*0} \to D^0\pi^0)} = 0.572 \pm 0.057 \pm 0.081$[5].

To be more closely connected with QCD dynamics, we have also calculated these form factors and decay widths based on the BS equation with a QCD-motivated interquark potential[14]. In the instantaneous approximation the BS equation

$$(\not{p}_1 - m_1)\chi_P(q)(\not{p}_2 + m_2) = \frac{i}{2\pi} \int d^4k \; G(\vec{P}, \vec{q} - \vec{k})\chi_P(k), \tag{22}$$

where $G(\vec{P}, \vec{q} - \vec{k})$ represents an "instantaneous" interquark potential in momentum space, can be reduced to the following equation for the three dimensional BS wave function

$$\Phi_{\vec{P}}(\vec{q}) = \int dq^0 \chi_P(q^0, \vec{q}), \tag{23}$$

$$(P^0 - E_1 - E_2)\Phi_{\vec{P}}(\vec{q}) = \Lambda_+^1 \gamma^0 \int d^3k \; G(\vec{P}, \vec{q} - \vec{k})\Phi_{\vec{P}}(\vec{k})\gamma^0 \Lambda_-^2, \tag{24}$$

where the contribution of negative energy projectors (i.e. the pair terms) are neglected.

The interquark potential is described by a long-ranged linear confining potential (Lorentz scalar V_S) plus a short-ranged one gluon exchange potential (Lorentz vector V_V), i.e.

$$\begin{aligned} V(r) &= V_S(r) + \gamma_\mu \otimes \gamma^\mu V_V(r), \\ V_S(r) &= \lambda r \frac{(1 - e^{-\alpha r})}{\alpha r}, \\ V_V(r) &= -\frac{4}{3} \frac{\alpha_s(r)}{r} e^{-\alpha r}, \end{aligned} \tag{25}$$

where the introduction of the factor $e^{-\alpha r}$ is to regulate the infrared (IR) divergence and also to incorporate the color screening effects of the dynamical light quark pairs on the $Q\bar{Q}$ potential. It is clear that when $\alpha r \ll 1$ the potentials given here become

identical with the standard linear plus Coulomb potential. In momentum space the potentials are

$$
\begin{aligned}
G(\vec{p}) &= G_S(\vec{p}) + \gamma_\mu \otimes \gamma^\mu G_V(\vec{p}), \\
G_S(\vec{p}) &= -\frac{\lambda}{\alpha}\delta^3(\vec{p}) + \frac{\lambda}{\pi^2}\frac{1}{(\vec{p}^{\,2}+\alpha^2)^2}, \\
G_V(\vec{p}) &= -\frac{2}{3\pi^2}\frac{\alpha_s(\vec{p})}{\vec{p}^{\,2}+\alpha^2},
\end{aligned}
\tag{26}
$$

where $\alpha_s(\vec{p})$ is the well known running coupling constant and is assumed to become a constant of $O(1)$ as $\vec{p}^{\,2} \to 0$

$$
\alpha_s(\vec{p}) = \frac{12\pi}{27}\frac{1}{\ln(a + \frac{\vec{p}^{\,2}}{\Lambda^2_{QCD}})}. \tag{27}
$$

The constants λ, α, a and Λ_{QCD} are the parameters that characterize the potential. In the computation we will use $\lambda = 0.18GeV^2,\ \ \alpha = 0.06GeV,\ \ a = e = 2.7183,\ \ \Lambda_{QCD} = 0.15GeV$.

Substituting (4), (5), and (26), (27) into the reduced BS equation (24), with quark masses $m_u = m_d = 200 - 350\ MeV,\ m_c = 1500\ MeV$, we can solve for the scalar wave functions ϕ and f of the 0^- and 1^- mesons respectively, which are different due to the color magnetic force induced by one gluon exchange potential. We then substitute the obtained ϕ and f into (11) and (17) to calculate the form factors. The results are also shown in Table 1.

The form factors are more suppressed for the wave functions by solving BS equation than for the Gaussian wave functions. This is mainly because, due to the relativistic correction of one gluon exchange potential (like the Breit-Fermi Hamiltonian) in the BS equation the spatial wave function of the 0^- meson has a larger $\langle \vec{q}^{\,2} \rangle$ (caused by an attractive spin-spin force between the quark and antiquark), and the 1^- meson wave function becomes different from the 0^- meson wave function, and therefore the overlap integral of wave functions between D^* and D is reduced. From Table 1 we see that although the decay widths are predicted to be somewhat different in the two models, the obtained values for g_A are close to each other. E.g., for $m_u = 300MeV$, $g_A = 0.65 - 0.67$; for $m_u = 200MeV$, $g_A = 0.81 - 0.83$. Nevertheless, we prefer the results obtained from BS equation with QCD-motivated potentials.

In the heavy quark limit ($m_Q \to \infty$) for the hadronic decay $D^* \to D\pi$ the heavy meson-pion strong coupling $G_{D^*D\pi}$, defined by

$$
\langle D^0(p_f)\pi^+(p_\pi) \mid D^{*+}(p_i,\epsilon)\rangle = G_{D^*D\pi}p^\mu_\pi \epsilon_\mu, \tag{28}
$$

may be written in a more convenient form as

$$G_{D^*D\pi} = \frac{2M_{D^*}}{F_\pi} g. \tag{29}$$

Comparing (28),(29), with (19) it is easy to find

$$g = g_A \xi_{A1}. \tag{30}$$

In the large N_C limit (g_A=1) and the nonrelativistic limit (ξ_{A1}=1) we would expect

$$g = 1. \tag{31}$$

However, based on the relativistic description for heavy meson decays, with a typical value of the constituent quark mass $m_u = 350(300)MeV$ our BS model gives

$$\xi_{A1} = 0.67(0.64), \quad g_A = 0.60(0.65), \quad g = 0.40(0.42). \tag{32}$$

We see that our estimated values of g_A are significantly smaller than (2): $g_A = 0.8 - 0.9$, which is expected in the chiral lagrangian approach. This is similar to the result of ref.[12]. Moreover, our value for the meson strong coupling $g \approx 0.4$ is also much smaller than (31): $g = 1$.

However, our result for the quark axial vector coupling g_A is sensitive to the value of the constituent quark mass. We see from Table 1 that if $m_u = 200MeV$ then the BS model calculation would give $\xi_{A1} = 0.58$, $g_A = 0.81$. In this case, although the light quark inside the heavy meson becomes even more relativistic, the $M1$ transition $D^{*0} \to D^0\gamma$ width gets enhanced due to a larger Dirac moment of the even lighter quark. Consequently, with the $CLEO$ ratio (3) the strong decay width of D^{*0} meson and the value of g_A will be increased. This possibility is of course not excluded and worth further investigating. In this connection it might be interesting to notice that $m_u \approx 200\ MeV$ is also suggested in some relativistic quark models (see, e.g., ref.[4]).

As for the meson strong coupling g, with $m_u = 200MeV$ we get $g = 0.47$, which is still much smaller than $g = 1$. In fact, our BS model calculations show that with a wide range of the constituent quark mass, say, $m_u = 200 - 350MeV$, the obtained meson strong coupling is $g = 0.47 - 0.40$. The Gaussian wave functions give $g = 0.53 - 0.45$ for $m_u = 200 - 300MeV$ (see Table 1). All these results suggest

$$g = 0.4 - 0.6, \tag{33}$$

which is much smaller than $g = 1$ expected in the large N_C and nonrelativistic limit, and than $g = 0.8 - 1$ suggested e.g. in refs.[8,19], but is consistent with the result $g \approx 0.6$ in a heavy hadron chiral model[9]. It is interesting to note that some recent QCD sum rule analyses favor an even smaller value $g = 0.2 - 0.4$[17]. The study of semileptonic decay $D \to \pi l \nu_l$ using a chiral effective theory in the heavy quark limit

also favors $g = 0.4$[18]. Moreover, $g = \frac{1}{3}$ is suggested by some calculation for the relativistic effects[13]. With (29) and (33) we obtain the following estimate for the effective $D^* D\pi$ coupling and $B^* B\pi$ coupling

$$G_{D^* D\pi} = 12 - 18, \qquad G_{B^* B\pi} = 32 - 48. \tag{34}$$

Finally, one might ask whether the nonrelativistic treatment for the light quarks is still tenable for heavy meson decays. In the nonrelativistic limit we would have $\xi_u = \xi_c = \xi_{A1} = 1$, then from (15) and the $CLEO$ ratio (3) with typical quark masses $m_c = 1500\ MeV$, $m_u = m_d = 330\ MeV$ we would get $\Gamma_{tot}(D^{*0}) = 107\ KeV$, $\Gamma_{tot}(D^{*+}) = 151 KeV$, which already exceeds the observed upper bound $\Gamma_{tot}(D^{*+}) \langle 131\ KeV$[6]. This seems to rule out the possibility that the light constituent quark inside the heavy meson can be treated as nonrelativistic, indicating that a relativistic description for the light quarks in heavy meson decays is necessary. It is amusing that by above nonrelativistic treatment we would get

$$g_A = 0.72, \tag{35}$$

which is almost in coincidence with the nonrelativistic value $g_A = 0.75$ obtained from the nucleon β-decay.

To sum up, we have estimated the form factors and decay rates for the $D^* \to D\gamma$ and $D^* \to D\pi$ decays, based on a relativistic constituent quark model with QCD inspired interquark potentials. Relativistic effects are found to be substantial on the form factors. Using the $CLEO$ experimental value $R^0_\gamma \equiv \frac{\Gamma(D^{*0} \to D^0 \gamma)}{\Gamma(D^{*0} \to D^0 \pi^0)} = 0.572 \pm 0.057 \pm 0.081$ as input we find that the quark axial vector coupling might be consistent with $g_A = 0.8 - 0.9$, which is expected in the chiral lagrangian approach. However, this can only be achieved with a smaller light quark mass, say, $m_u = m_d = (200 - 220) MeV$, which give larger $M1$ transition widths accordingly. With the typical value $m_u = m_d = 300 - 350 MeV$, our obtained value for g_A is smaller, say $0.67 - 0.60$. It is therefore interesting to note that the g_A issue might be possibly related to the value of the constituent quark mass of light quarks in the heavy meson decays. As for the heavy meson-pion strong coupling g, with a wide range of the constituent quark mass $m_u = m_d = (200 - 350) MeV$, we obtain $g = 0.6 - 0.4$, which is much smaller than $g = 1$ expected in the large N_C and nonrelativistic limit.

	Gaussian		*BS*		
	$m_u = 200$ (MeV)	$m_u = 300$ (MeV)	$m_u = 200$ (MeV)	$m_u = 300$ (MeV)	$m_u = 350$ (MeV)
ξ_u	0.42	0.51	0.36	0.46	0.51
ξ_c	0.97	0.95	0.90	0.87	0.86
ξ_{A1}	0.64	0.70	0.58	0.64	0.67
g_A	0.83	0.65	0.81	0.65	0.60
g	0.53	0.45	0.47	0.42	0.40
$\Gamma(D^{*0} \to D^0\gamma)$	$21(KeV)$	$15(KeV)$	$16(KeV)$	$13(KeV)$	$12(KeV)$
$\Gamma(D^{*0} \to D^0\pi^0)$	37	26	29	22	21
$\Gamma_{tot}(D^{*0})$	58	41	45	35	33
$\Gamma(D^{*+} \to D^+\gamma)$	0.42	0.12	0.24	0.10	0.07
$\Gamma(D^{*+} \to D^+\pi^0)$	26	18	20	15	14
$\Gamma(D^{*+} \to D^0\pi^+)$	55	40	43	33	31
$\Gamma_{tot}(D^{*+})$	81	58	63	48	46

Table 1: Form factors and decay widths for $D^* \to D\gamma$ and $D^* \to D\pi$ decays. Predicted values are given with (1) the Gaussian wave functions and (2) the wave functions by solving BS equation. Here $m_c = 1500MeV, m_u = m_d = 200,\ 300,\ 350\ MeV$ are assumed. The CLEO experimental value of $R_\gamma^0 \equiv \frac{\Gamma(D^{*0}\to D^0\gamma)}{\Gamma(D^{*0}\to D^0\pi^0)} = 0.572 \pm 0.057 \pm 0.081$ is used as input.

This work was supported in part by the National Natural Science Foundation of China, and the State Education Commission of China. After this work was submitted for publication we received the Bari preprints (refs.[13,17]), and one of us (K.T.C) would like to thank Dr.De Fazio and Prof.Nardulli for the useful communication.

References

[1] S. Weinberg, Phys. Rev. Lett. **65** (1990) 1181.

[2] S. Weinberg, Phys. Rev. Lett. **67** (1991) 3473;
D.A. Dicus et al. , Phys. Lett. B **284** (1992)384.

[3] S. Peris, Phys. Lett. B **268** (1991)415;
S. Peris, Phys. Rev. D **46** (1992)1202.

[4] C. Hayne and N. Isgur, Phys. Rev. D**25** (1982)1944.

[5] CLEO Collaboration, F.Butler et al.,Phys. Rev. Lett. **69** (1992) 2041.

[6] Particle Data Group, L. Montanet et al., Phys. Rev. D **50** (1994)1171.

[7] M.B. Wise, Phys. Rev. D**45** (1992) R2188.

[8] T.M. Yan et al., Phys. Rev. D**46** (1992)1148.

[9] P. Cho and H. Georgi, Phys.Lett.B **296** (1992) 408; B **300** (1993)410 (E).

[10] J.F. Amundson et al., Phys.Lett.B **296** (1992) 415.

[11] H.Y. Cheng et al., Phys.Rev.D**47** (1993) 527.

[12] Y.B. Dai, C.S. Huang, and J.H. Ying, Z. Phys.C **60** (1993) 527.

[13] P. Colangelo, F.De Fazio, and G. Nardulli, Bari University preprint BARI-TH/94-180.

[14] K.T. Chao, talk given at the ITP Workshop on Hadron Transition Matrix Elements, Beijing, November 1990, unpublished;
J.H. Liu, Ph. D Thesis, Peking University (May 1993), unpublished;
J. Tang, J.H. Liu and K.T. Chao, Phys.Rev.D, to be published.

[15] H.Georgi and A.Manohar, Nucl.Phys.B**234**(1984)189.

[16] N.Isgur, D.Scora, B.Grinstein and M.B.Wise, Phys. Rev. D**39**(1989)799.

[17] P. Colangelo et al., Bari University preprint BARI-TH/94-171.

[18] P. Cogangelo, F.De Fazio, and G. Nardulli, Phys. Lett. B**316** (1993)555;
R. Casalbuoni et al., Phys. Lett. B **299**(1993)139.

[19] N. Isgur and M. B. Wise, Phys. Rev. D **41**(1990)151.